

Two Loop Neutrino Model with Dark Matter and Leptogenesis

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Abstract

We study a two-loop induced radiative neutrino model at TeV scale with global $U(1)$ symmetry, in which we analyze dark matter and resonant leptogenesis. The model includes two kinds of dark matter candidates. We discuss what kind of dark matter can satisfy the observed relic density as well as the current direct detection bound, and be simultaneously compatible with the leptogenesis. We also discuss whether our resonant leptogenesis can be differentiated from the other scenarios at TeV scale or not.

Keywords: Radiative seesaw, Global $U(1)$ symmetry, Dark matter, Resonant leptogenesis

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I. INTRODUCTION

After the discovery of the Higgs boson at the LHC, the Standard Model (SM) has been established well. However the SM still has to be extended in order to explain the existence of dark matter (DM), the small neutrino masses, the baryon asymmetry in the universe and so on. Radiative seesaw scenarios are renowned as one of the economical models which simultaneously explain the existence of DM and the neutrino masses. Since the DM candidate is necessary to generate the neutrino masses in this kind of models, physics between DM and neutrinos is strongly correlated in this simple framework. For example, couplings and mass scale of DM are related with the neutrino mass scale. In addition, since this kind of models can naturally include a new particle with TeV scale mass, radiative seesaw scenarios have good testability in near future experiments. Along this line of idea, a vast literature has recently arisen in Ref. [1–72].

On the other hand, explaining the observed baryon asymmetry in the universe via leptogenesis is one of the challenging issues in the framework of radiative seesaw models with right-handed neutrinos, since couplings related to the source of the leptogenesis are expected to be $\mathcal{O}(1)$ due to the requirement of the neutrino masses. It causes the strong washout of the generated baryon asymmetry. In order to avoid this problem, we have to take hierarchical couplings with highly degenerated masses between the source and the mediated fields.¹ For example, generating the baryon asymmetry via resonant leptogenesis has been discussed based on the Ma model [73].

In this paper, we study a two-loop induced radiative neutrino model at TeV scale with a global $U(1)$ symmetry, in which we analyze DM and resonant leptogenesis simultaneously. In this model, we have a scalar or a fermion DM candidate. We discuss which kind of DM candidate can satisfy the observed relic density as well as the current direct detection bound, and can also be compatible with leptogenesis. Since our model has two sources of leptogenesis, we also show different points of our resonant leptogenesis from the other scenarios such as the Ma model at TeV scale [73].

This paper is organized as follows. In Sec. II, we show our model including field contents and their global $U(1)$ charges, Higgs potential, and neutrino masses. In Sec. III, DM properties including relic density and current limit by direct detection experiments are discussed. In Sec. IV, we analyze resonant leptogenesis. Summary and conclusions are given in Sec. V.

¹ Such couplings can be achieved by making use of the experimental fact that one of three neutrino masses can be negligible. To get the sufficient baryon asymmetry via thermal leptogenesis in the radiative seesaw framework, resonant leptogenesis would be only the solution that requires the mass degeneracy between the source and the mediated fields.

II. THE MODEL

A. Model setup

	L_{Li}	e_{Ri}	$F_{L/Rj}$	N_{Rj}	X_{Rj}	Φ	η	χ^0	$\chi^{0'}$	Σ
$(SU(2)_L, U(1)_Y)$	$(\mathbf{2}, -1/2)$	$(\mathbf{1}, -1)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$	$(\mathbf{2}, 1/2)$	$(\mathbf{2}, 1/2)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$
$U(1)$	$-x/2$	$-x/2$	x	$3x/2$	$-3x/2$	0	$3x/2$	$-x/2$	$-5x/2$	x
Accidental \mathbb{Z}_2	+	+	-	+	+	+	-	-	-	+

TABLE I: Field contents and charge assignments of $SU(2)_L \times U(1)_Y \times U(1)$, where indices $i = 1 - 3$ and $j = 1, 2, (3)$ represent the generation.

As shown in Tab. I, we introduce two (or three) gauge singlet vector-like fermions $F_{L/R}$, and gauge singlet Majorana fermions N_R , and X_R as new fermions. The number of these particles should be more than two in order to obtain at least two non-zero neutrino mass eigenvalues. We also introduce an inert $SU(2)_L$ doublet scalar η , two neutral inert singlet scalars $(\chi^0, \chi^{0'})$, and a neutral singlet scalar Σ as new scalars. We assume that only the Higgs doublet field in the SM Φ and the new singlet scalar Σ have vacuum expectation values (VEVs), which are symbolized by $\langle \Phi \rangle = v/\sqrt{2}$ and $\langle \Sigma \rangle = v'/\sqrt{2}$ respectively.² We impose a global $U(1)$ symmetry, under which Φ does not have a charge in order not to couple to the Goldstone boson (GB) [37]. The global $U(1)$ charge $x \neq 0$ is in principle an arbitrary, and the field assignments play a crucial role in realizing our neutrino masses at two-loop level. If the $U(1)$ charge x is fixed to be $x = 2$, we can identify this $U(1)$ symmetry as a global $B - L$ symmetry. Hereafter we assume this global $U(1)$ symmetry to be a kind of $U(1)_{B-L}$ symmetry. Note that while the new fermions are added as vector-like and do not contribute to anomalies, this model is anomalous since the three chiral fermions with $B - L = -1$ corresponding to the right-handed neutrinos are not introduced. If one would like to have an anomaly free model, the anomalies can be cancelled by introducing some pairs of new heavy vector-like fermions [74, 75]. However this is beyond the scope of this paper. This model has an accidental \mathbb{Z}_2 symmetry which can assure the DM stability, and the accidental \mathbb{Z}_2 assignments are shown in Tab. I.

The renormalizable Lagrangian for Yukawa sector, mass term, and scalar potential under the

² The scale of v' should be larger than $v' \sim 10^7$ GeV for successful leptogenesis as we will see later. Otherwise the annihilation channel $N_1 N_1, X_1 X_1 \rightarrow GG$ whose reaction rate is determined by v' does not satisfy the out-of-equilibrium condition at $T \sim \mathcal{O}(10)$ TeV where T is the temperature of the universe. Thus the baryon asymmetry would be washed out due to this process.

charge assignments are given by

$$\begin{aligned}\mathcal{L}_Y = & -y_\ell \bar{L}_L \Phi e_R - y_\eta \bar{L}_L \eta^\dagger F_R - y_{N\chi} \bar{F}_L N_R \chi^0 - y_{N\chi'} \bar{F}_L X_R \chi^{0\dagger} \\ & - y'_{N\chi'} \bar{F}_R^c N_R \chi^{0'} - y'_{N\chi} \bar{F}_R^c X_R \chi^{0\dagger} - M_{NX} \bar{N}_R^c X_R - M_F \bar{F}_L F_R + \text{H.c.},\end{aligned}\quad (\text{II.1})$$

$$\begin{aligned}\mathcal{V} = & m_\Phi^2 |\Phi|^2 + m_\eta^2 |\eta|^2 + m_\chi^2 |\chi^0|^2 + m_{\chi'}^2 |\chi^{0'}|^2 + m_\Sigma^2 |\Sigma|^2 \\ & + \left[\lambda (\Phi^\dagger \eta) \chi^0 \Sigma^\dagger + \lambda' (\Phi^\dagger \eta) \chi^{0'} \Sigma + \frac{\lambda''}{2} (\chi^{0\dagger} \chi^{0'}) \Sigma^{\dagger 2} + \frac{\mu_\chi}{2} (\chi^0)^2 \Sigma + \text{H.c.} \right] \\ & + \frac{\lambda_\Phi}{4} |\Phi|^4 + \frac{\lambda_\eta}{4} |\eta|^4 + \frac{\lambda_\chi}{4} |\chi^0|^4 + \frac{\lambda_{\chi'}}{4} |\chi^{0'}|^4 + \frac{\lambda_\Sigma}{4} |\Sigma|^4 + \lambda_{\Phi\eta} |\Phi|^2 |\eta|^2 + \lambda'_{\Phi\eta} (\Phi^\dagger \eta) (\eta^\dagger \Phi) \\ & + \lambda_{\Phi\chi} |\Phi|^2 |\chi^0|^2 + \lambda_{\Phi\chi'} |\Phi|^2 |\chi^{0'}|^2 + \lambda_{\Phi\Sigma} |\Phi|^2 |\Sigma|^2 + \lambda_{\eta\chi} |\eta|^2 |\chi^0|^2 + \lambda_{\eta\chi'} |\eta|^2 |\chi^{0'}|^2 + \lambda_{\eta\Sigma} |\eta|^2 |\Sigma|^2 \\ & + \lambda_{\chi\chi'} |\chi^0|^2 |\chi^{0'}|^2 + \lambda_{\chi\Sigma} |\chi^0|^2 |\Sigma|^2 + \lambda_{\chi'\Sigma} |\chi^{0'}|^2 |\Sigma|^2,\end{aligned}\quad (\text{II.2})$$

where the first term in \mathcal{L}_Y generates the SM charged lepton masses, and we assume the couplings λ , λ' , λ'' and μ_χ in the scalar potential to be real for simplicity. As we will see below, these couplings become important for neutrino mass generation.

After the symmetry breaking, the scalar fields can be parametrized by

$$\Phi = \begin{pmatrix} w^+ \\ \frac{1}{\sqrt{2}}(v + \phi + iz) \end{pmatrix}, \quad \eta = \begin{pmatrix} \eta^+ \\ \frac{1}{\sqrt{2}}(\eta_R + i\eta_I) \end{pmatrix}, \quad \Sigma = \frac{v' + \sigma}{\sqrt{2}} e^{iG/v'}.\quad (\text{II.3})$$

where $v \approx 246$ GeV is the VEV of the SM Higgs doublet, and w^\pm and z are respectively the GBs which are absorbed by the longitudinal components of the W and Z bosons. Inserting the tadpole conditions, the resulting mass matrix of the CP even scalar (ϕ, σ) is given by

$$m^2(\phi, \sigma) = \begin{pmatrix} \lambda_\Phi v^2 & 2\lambda_{\Phi\Sigma} v v' \\ 2\lambda_{\Phi\Sigma} v v' & \lambda_\Sigma v'^2 \end{pmatrix} = \begin{pmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} m_h^2 & 0 \\ 0 & m_H^2 \end{pmatrix} \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix}, \quad (\text{II.4})$$

where h is the SM-like Higgs boson and H is an additional CP-even Higgs mass eigenstate. The gauge eigenstates ϕ and σ are rewritten in terms of the mass eigenstates h and H as

$$\begin{pmatrix} \phi \\ \sigma \end{pmatrix} = \begin{pmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} h \\ H \end{pmatrix}, \quad \text{with} \quad \sin 2\alpha = \frac{4\lambda_{\Phi\Sigma} v v'}{m_H^2 - m_h^2}.\quad (\text{II.5})$$

The GB G in Eq. (II.3) appears due to the spontaneous symmetry breaking of the global $U(1)$ symmetry. The couplings between the GB and the particles with non-trivial global $U(1)$ charges are given by $J^\mu \partial_\mu G/v'$ through the global $U(1)$ current J^μ . As one can see, the coupling is suppressed by the VEV v' .

The mass matrices of the CP even and CP odd states of the inert scalar bosons $(\eta, \chi^0, \chi^{0'})_{R/I}$

are respectively given by

$$M_R^2 = \begin{pmatrix} m_\eta^2 + \frac{(\lambda_{\Phi\eta} + \lambda'_{\Phi\eta})v^2 + \lambda_{\eta\Sigma}v'^2}{2} & \lambda vv'/2 & \lambda' vv'/2 \\ \lambda vv'/2 & m_\chi^2 + \frac{\sqrt{2}\mu_\chi v' + \lambda_{\Phi\chi}v^2 + \lambda_{\chi\Sigma}v'^2}{2} & \lambda'' v'^2/4 \\ \lambda' vv'/2 & \lambda'' v'^2/4 & m_{\chi'}^2 + \frac{\lambda_{\Phi\chi'}v^2 + \lambda_{\chi'\Sigma}v'^2}{2} \end{pmatrix}, \quad (\text{II.6})$$

$$M_I^2 = \begin{pmatrix} m_\eta^2 + \frac{(\lambda_{\Phi\eta} + \lambda'_{\Phi\eta})v^2 + \lambda_{\eta\Sigma}v'^2}{2} & -\lambda vv'/2 & -\lambda' vv'/2 \\ -\lambda vv'/2 & m_\chi^2 + \frac{-\sqrt{2}\mu_\chi v' + \lambda_{\Phi\chi}v^2 + \lambda_{\chi\Sigma}v'^2}{2} & -\lambda'' v'^2/4 \\ -\lambda' vv'/2 & -\lambda'' v'^2/4 & m_{\chi'}^2 + \frac{\lambda_{\Phi\chi'}v^2 + \lambda_{\chi'\Sigma}v'^2}{2} \end{pmatrix}, \quad (\text{II.7})$$

where we define diagonal mass matrices $(M_d^2)_{R/I} \equiv (m_{R_1/I_1}^2, m_{R_2/I_2}^2, m_{R_3/I_3}^2)$, and their mixing matrices $O_{id}^{R/I}$, so that they satisfy $M_{R/I}^2 \equiv O_{id}^{R/I}(M_d^2)_{R/I}(O_{dj}^{R/I})^T$. Depending on the couplings, the lightest CP even or CP odd mass eigenstate with the mass m_{R_1} or m_{I_1} can be a DM candidate.³ The non-standard couplings between DM and the GB induced by the non-self-conjugate couplings λ , λ' , λ'' and μ_χ may be relevant to compute the DM relic density. This coupling can be written down as

$$\mathcal{V}_{\text{DM-DM-G-G}} = - \left(\frac{\lambda}{4} \frac{v}{v'} O_{11}^R O_{21}^R + \frac{\lambda'}{4} \frac{v}{v'} O_{11}^R O_{31}^R + \frac{\lambda''}{2} O_{21}^R O_{31}^R + \frac{\mu_\chi}{4\sqrt{2}v'} (O_{21}^R)^2 \right) \text{DM}^2 G^2, \quad (\text{II.8})$$

with the mixing matrix O^R . In addition, the couplings between the CP even Higgs bosons and the GB are also relevant to compute the DM relic density. These couplings come from the kinetic term of Σ and can be written as

$$\mathcal{L} \supset \left[\frac{-\sin \alpha h + \cos \alpha H}{v'} + \frac{(-\sin \alpha h + \cos \alpha H)^2}{2v'^2} \right] (\partial_\mu G) (\partial^\mu G). \quad (\text{II.9})$$

Finally the mass eigenvalue of the charged inert scalar η^\pm is given by

$$m_{\eta^\pm}^2 = m_\eta^2 + \frac{\lambda_{\Phi\eta}v^2 + \lambda_{\eta\Sigma}v'^2}{2}. \quad (\text{II.10})$$

In this model, the typical mass scale of these new exotic particles is assumed to be TeV scale. On the other hand, we should take $v' \sim 10^7$ GeV for successful leptogenesis. Therefore a certain degree of tuning among the relevant couplings cannot be avoided. More specifically, demanding that the diagonal elements of the mass matrix Eq. (II.4), (II.6) and (II.7) are TeV scale and off-diagonal elements are 10 GeV scale to obtain small mixings of the order of $O_{ij}^{R/I} \sim 10^{-2}$ ($i \neq j$), the order of magnitude of the couplings should roughly be $\lambda'' \sim 10^{-12}$, λ , λ' , $\lambda_{\Phi\Sigma} \sim 10^{-7}$ and $\lambda_{\eta\Sigma}$, $\lambda_{\chi\Sigma}$, $\lambda_{\chi'\Sigma}$, $\lambda_\Sigma \sim 10^{-8}$. Although this point may be a disadvantage of this model, it would be worth

³ As we will discuss later, the CP even state is identified as DM.

discussing such a new concrete model with a global $U(1)$ symmetry as an example model since all the phenomenology of the neutrino masses, the existence of DM and the baryon asymmetry of the universe are closely correlated.

B. Neutrino mass matrix

Due to renormalizability and the strong restriction of interactions via the global $U(1)$ symmetry in this model, neutrino masses are not generated neither tree level nor one-loop level. If the vector like fermion F has a Majorana mass term, neutrino masses would be generated at one-loop level (for example see Ref. [1]), however this is not our case. As a result, neutrino masses are induced at two-loop level, and we have three types of diagrams as shown in Fig. 1. The formula of the total neutrino mass matrix can be given by

$$\begin{aligned}
(m_\nu)_{\alpha\beta} = & - \sum_{i,j,k} \sum_{m,n} \frac{M_{F_i} M_{F_k}}{4M_{NX_j}} (y_\eta)_{\alpha i} \left[(y_{N\chi})_{ij} (y_{N\chi'})_{kj} + (y_{N\chi'})_{ij} (y_{N\chi})_{kj} \right] (y_\eta)_{\beta k} \\
& \times \left[I_{1(mn)}^{(ijk)} + I_{2R(mn)}^{(ijk)} + I_{3R(mn)}^{(ijk)} \right] \\
& - \sum_{i,j,k} \sum_{m,n} \frac{M_{F_i} M_{F_k}}{4M_{NX_j}} (y_\eta)_{\alpha i} \left[(y'_{N\chi})_{ij} (y'^*_{N\chi'})_{kj} + (y'^*_{N\chi'})_{ij} (y'_{N\chi})_{kj} \right] (y_\eta)_{\beta k} \\
& \times \left[I_{2L(mn)}^{(ijk)} + I_{3L(mn)}^{(ijk)} \right], \tag{II.11}
\end{aligned}$$

where $I_{1(mn)}^{(ijk)}$, a pair of $I_{2R(mn)}^{(ijk)}$ and $I_{2L(mn)}^{(ijk)}$, a pair of $I_{3R(mn)}^{(ijk)}$ and $I_{3L(mn)}^{(ijk)}$ are the dimensionless loop functions which come from the left, center and right diagrams in Fig. 1 respectively. These loop functions are defined by

$$I_{1(mn)}^{(ijk)} = O_{1213}^{R_m R_n} I_{1(R_m R_n)}^{(ijk)} - O_{1213}^{R_m I_n} I_{1(R_m I_n)}^{(ijk)} + O_{1213}^{I_m R_n} I_{1(I_m R_n)}^{(ijk)} - O_{1213}^{I_m I_n} I_{1(I_m I_n)}^{(ijk)}, \tag{II.12}$$

$$I_{2R(mn)}^{(ijk)} = O_{1312}^{R_m R_n} I_{2R(R_m R_n)}^{(ijk)} + O_{1312}^{R_m I_n} I_{2R(R_m I_n)}^{(ijk)} - O_{1312}^{I_m R_n} I_{2R(I_m R_n)}^{(ijk)} - O_{1312}^{I_m I_n} I_{2R(I_m I_n)}^{(ijk)}, \tag{II.13}$$

$$I_{2L(mn)}^{(ijk)} = O_{1213}^{R_m R_n} I_{2L(R_m R_n)}^{(ijk)} - O_{1213}^{R_m I_n} I_{2L(R_m I_n)}^{(ijk)} + O_{1213}^{I_m R_n} I_{2L(I_m R_n)}^{(ijk)} - O_{1213}^{I_m I_n} I_{2L(I_m I_n)}^{(ijk)}, \tag{II.14}$$

$$I_{3R(mn)}^{(ijk)} = O_{1123}^{R_m R_n} I_{3R(R_m R_n)}^{(ijk)} + O_{1123}^{R_m I_n} I_{3R(R_m I_n)}^{(ijk)} - O_{1123}^{I_m R_n} I_{3R(I_m R_n)}^{(ijk)} - O_{1123}^{I_m I_n} I_{3R(I_m I_n)}^{(ijk)}, \tag{II.15}$$

$$I_{3L(mn)}^{(ijk)} = O_{1123}^{R_m R_n} I_{3L(R_m R_n)}^{(ijk)} + O_{1123}^{R_m I_n} I_{3L(R_m I_n)}^{(ijk)} - O_{1123}^{I_m R_n} I_{3L(I_m R_n)}^{(ijk)} - O_{1123}^{I_m I_n} I_{3L(I_m I_n)}^{(ijk)}, \tag{II.16}$$

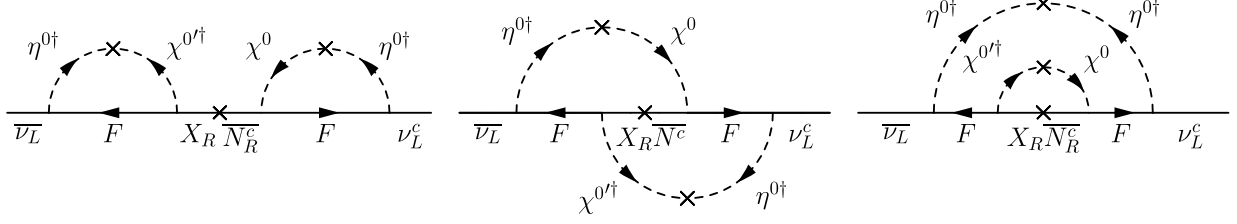


FIG. 1: Radiative generation of neutrino masses.

where $O_{abcd}^{R_m I_n} = O_{am}^R O_{bm}^R O_{cn}^I O_{dn}^I$ and

$$I_{1(R_m I_n)}^{(ijk)} = \frac{1}{(4\pi)^4} \left[I \left(\frac{m_{R_m}^2}{M_{F_i}^2} \right) I \left(\frac{m_{I_n}^2}{M_{F_k}^2} \right) + I \left(\frac{m_{R_m}^2}{M_{F_k}^2} \right) I \left(\frac{m_{I_n}^2}{M_{F_i}^2} \right) \right], \quad (\text{II.17})$$

$$I_{2R(R_m I_n)}^{(ijk)} = M_{N X_j}^2 \int \frac{d^4 \ell}{(2\pi)^4} \int \frac{d^4 q}{(2\pi)^4} \frac{1}{\ell^2 - M_{F_i}^2} \frac{1}{(\ell - q)^2 - M_{N X_j}^2} \frac{1}{q^2 - M_{F_k}^2} \frac{1}{\ell^2 - m_{R_m}^2} \frac{1}{q^2 - m_{I_n}^2}, \quad (\text{II.18})$$

$$I_{2L(R_m I_n)}^{(ijk)} = -\frac{M_{N X_j}^2}{4M_{F_i} M_{F_k}} \int \frac{d^4 \ell}{(2\pi)^4} \int \frac{d^4 q}{(2\pi)^4} \frac{\ell \cdot q}{\ell^2 - M_{F_i}^2} \frac{1}{(\ell - q)^2 - M_{N X_j}^2} \frac{1}{q^2 - M_{F_k}^2} \frac{1}{\ell^2 - m_{R_m}^2} \frac{1}{q^2 - m_{I_n}^2}, \quad (\text{II.19})$$

$$I_{3R(R_m I_n)}^{(ijk)} = M_{N X_j}^2 \int \frac{d^4 \ell}{(2\pi)^4} \int \frac{d^4 q}{(2\pi)^4} \frac{1}{\ell^2 - M_{F_i}^2} \frac{1}{q^2 - M_{N X_j}^2} \frac{1}{\ell^2 - M_{F_k}^2} \frac{1}{\ell^2 - m_{R_m}^2} \frac{1}{(\ell - q)^2 - m_{I_n}^2}, \quad (\text{II.20})$$

$$I_{3L(R_m I_n)}^{(ijk)} = -\frac{M_{N X_j}^2}{M_{F_i} M_{F_k}} \int \frac{d^4 \ell}{(2\pi)^4} \int \frac{d^4 q}{(2\pi)^4} \frac{\ell^2}{\ell^2 - M_{F_i}^2} \frac{1}{q^2 - M_{N X_j}^2} \frac{1}{\ell^2 - M_{F_k}^2} \frac{1}{\ell^2 - m_{R_m}^2} \frac{1}{(\ell - q)^2 - m_{I_n}^2}, \quad (\text{II.21})$$

with $I(x) = x \log x / (1 - x)$. Note that in the derivation of the above formula, the CP phases except the Yukawa couplings are neglected. The contribution of the left diagram can be understood as linear seesaw like formula by splitting the diagram into two Dirac masses induced at one-loop level. For the center and right diagrams, there are two kinds of contributions coming from right and left chiralities of the internal fermions. In other words, these two contributions to the neutrino masses come from the masses or momenta of the F propagators in the loop respectively. The neutrino mass generation can be understood as follows. Due to the global $U(1)$ symmetry breaking by the VEV of Σ , the mixing between η , χ^0 and $\chi^{0'}$ occurs. Then since the global $U(1)$ symmetry is correlated with the lepton number conservation, the $U(1)$ symmetry breaking implies breaking of the lepton number. Thus the neutrino Majorana mass term is generated after the $U(1)$ symmetry breaking.

The neutrino mass matrix computed above can be diagonalized by the Pontecorvo-Maki-

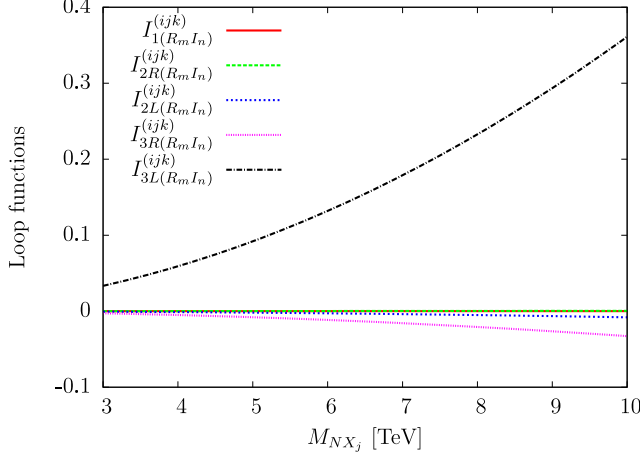


FIG. 2: Numerical calculation of the loop functions where the other masses are fixed as $M_{F_i} = M_{F_k} = 1.5$ TeV and $m_{R_m} = m_{I_n} = 1.2$ TeV which are typical sample points to discuss DM and leptogenesis as we will see below.

Nakagawa-Sakata matrix U_{PMNS} [76]; $U_{\text{PMNS}}^T(m_\nu)U_{\text{PMNS}} = \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$. The neutrino masses and their mixing angles are measured by experiments [77], and these values depend on normal or inverted mass hierarchy. In our model, the order of magnitude of the neutrino masses can roughly be estimated as $m_\nu \sim y_\eta^2 Y^2 O_{\text{mix}}^4 I_{\text{loop}} v'$ where Y is the dominant Yukawa coupling in $y_{N\chi}$, $y_{N\chi'}$, $y'_{N\chi}$, $y'_{N\chi'}$, O_{mix} represents the mixing matrix of O^R , O^I , and I_{loop} is the loop function. Thus one can find that the order of $y_\eta^2 Y^2 \sim 10^{-8}$ is required to obtain the experimental value $m_\nu \sim 0.1$ eV with the typical assumed mixing angle $O_{\text{mix}} \sim 10^{-2}$, $I_{\text{loop}} \sim 0.1$ and $v' \sim 10^7$ GeV. Note that the neutrino mass matrix should be proportional to the VEV v' since v' is the origin of the lepton number violation.

We numerically compute the loop functions with the public code SecDec [78] in order to evaluate the neutrino masses more precisely in this model. Here one should note that the loop functions $I_{3R(R_m I_n)}^{(ijk)}$ and $I_{3L(R_m I_n)}^{(ijk)}$ include a divergence. This is obvious from the definition of the loop functions in Eq. (II.20) and (II.21). However the divergent terms eventually cancel out with each other as follows. The loop functions can be regularized with dimensional regularization and expanded around dimension $d = 4$ which can be done within SecDec. With a brief evaluation, one can see that the terms including divergences are independent on at least either of m or n which is the index of the scalar mass eigenvalues. Thus the loop function including a divergence $I_{3R/L(R_m I_n)}^{(ijk)UV}$ can be written as $I_{3R/L(R_m I_n)}^{(ijk)UV} = I_{3R/L}^{(ijk)UV1} + I_{3R/L(R_m)}^{(ijk)UV2} + I_{3R/L(I_n)}^{(ijk)UV3}$. Taking into account this fact and the orthogonality of the mixing matrices O^R and O^I which means $\sum_m O_{im}^R O_{jm}^R = \sum_m O_{im}^I O_{jm}^I = \delta_{ij}$,

one can see that the first two terms $I_{3R/L}^{(ijk)UV1}$ and $I_{3R/L(R_m)}^{(ijk)UV2}$ vanish after taking the summation over n . Moreover the remaining third divergent term also cancels after all the relevant terms are summed in Eq. (II.15) and (II.16) as

$$\begin{aligned} \sum_{m,n} I_{3R/L(mn)}^{(ijk)UV} &= \sum_n \left[O_{2n}^R O_{3n}^R I_{3R/L(R_n)}^{(ijk)UV3} + O_{2n}^I O_{3n}^I I_{3R/L(I_n)}^{(ijk)UV3} - O_{2n}^R O_{3n}^R I_{3R/L(R_n)}^{(ijk)UV3} - O_{2n}^I O_{3n}^I I_{3R/L(I_n)}^{(ijk)UV3} \right] \\ &= 0. \end{aligned} \quad (\text{II.22})$$

Thus the divergent terms do not contribute to the neutrino masses. The numerical value of the loop functions are almost fixed by the maximum mass in M_{F_i} , M_{NX_j} , M_{F_k} , m_{R_m} and m_{I_n} , and the result obtained by using SecDec is shown in Fig. 2.⁴ The numerical calculation shows that the loop function $I_{3L(R_m I_n)}^{(ijk)}$ coming from the right diagram in Fig. 1 gives a dominant contribution to the neutrino masses.

C. LFV processes

We should take into account lepton flavor violations (LFVs) such as $\mu \rightarrow e\gamma$, which typically provide strong constraints on radiative neutrino mass models. In our case, such processes arise through only the y_η term, and analyses are very similar with the case of the Ma model [79], and the Yukawa couplings $y_{N\chi}$, $y_{N\chi'}$, $y'_{N\chi}$, $y'_{N\chi'}$ are not constrained by the LFV processes at least at one-loop level. Among the LFV processes, we focus on the one-loop induced $\mu \rightarrow e\gamma$ that gives the most stringent constraint on y_η and the mediating particles F and η^+ . The resulting formula for $\mu \rightarrow e\gamma$ and its experimental bound [80] are given by

$$\text{Br}(\mu \rightarrow e\gamma) = \frac{3\alpha_{\text{em}}}{64\pi G_F^2 m_{\eta^\pm}^4} \left| \sum_{i=1}^3 (y_\eta^*)_{i1} (y_\eta)_{i2} F_2(\xi_i) \right|^2 \leq 5.7 \times 10^{-13}, \quad (\text{II.23})$$

$$\text{with } F_2(\xi_i) = \frac{1 - 6\xi_i + 2\xi_i^3 + 3\xi_i^2 - 6\xi_i^2 \ln \xi_i}{6(1 - \xi_i)^4}, \quad (\text{II.24})$$

where $G_F = 1.17 \times 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant, $\xi_i = M_{F_i}^2/m_{\eta^\pm}^2$ and $\alpha_{\text{em}} = e^2/(4\pi) \approx 1/137$ is the electromagnetic fine structure constant. The simplest way to avoid this constraint is to assume y_η to be diagonal, since its formula is proportional to the off-diagonal elements of y_η as one can see in Eq. (II.23). In this case we expect the neutrino mixings can be derived through the other Yukawa couplings $y_{N\chi}$, $y_{N\chi'}$, $y'_{N\chi}$, $y'_{N\chi'}$. Otherwise the parameters are constrained as $y_\eta \lesssim 0.01$ and $m_{\eta^\pm} \gtrsim 200 \text{ GeV}$. For instance, with the values $y_\eta = 0.01$, $m_{\eta^\pm} = 200 \text{ GeV}$ and $F_2(\xi_i) = 1/6$, the maximum branching ratio is found to be $\text{Br}(\mu \rightarrow e\gamma) \approx 1.4 \times 10^{-13}$.

⁴ At most 1% error is included in the numerical calculation.

Since we take $y_\eta \lesssim \mathcal{O}(0.01)$ for the LFV constraint and $\mathcal{O}(1)$ TeV of the new particle masses in the loop, the other Yukawa couplings $y_{N\chi}$, $y_{N\chi'}$, $y'_{N\chi}$, $y'_{N\chi'}$ should be roughly larger than 10^{-2} in order to obtain the scale of the observed neutrino mass $m_\nu \sim 0.1$ eV assuming $O_{\text{mix}} \sim 10^{-2}$ as discussed in the previous section.

D. Goldstone Boson

Here we mention some issues on the GB. Due to the direct consequence of our global $U(1)$ symmetry, the GB remains as a physical state, which could be constrained by some experiments. In our case, the constraint comes from the invisible decay of the SM Higgs boson, and its decay width can be computed with the coupling given in Eq. (II.9) to be $\Gamma(h \rightarrow GG) = m_h^3 \sin^2 \alpha / (32\pi v'^2)$. This decay width should be smaller than 1.2 MeV at 95% confidential level [81], and thus we get the constraint

$$\left(\frac{\sin \alpha}{0.1}\right) \left(\frac{1 \text{ TeV}}{v'}\right) \lesssim 2.5. \quad (\text{II.25})$$

Moreover, $\sin \alpha$ itself is constrained by the latest LHC searches by ATLAS and CMS to be (conservatively) $\sin \alpha \lesssim 0.2$ [82]. Therefore the above constraint is translated to the constraint on the VEV as $v' \gtrsim 800$ GeV. However since we take $v' \sim 10^7$ GeV for successful leptogenesis, this bound is easily satisfied in our case.

Another bound comes from the Supernova 1987A observations and simulations, which tell us the following relation [83]:

$$|\lambda_{\Phi\Sigma}| \lesssim 0.011 \left(\frac{m_H}{500 \text{ MeV}}\right)^2. \quad (\text{II.26})$$

This bound also does not affect to our model seriously, since both m_H and $\lambda_{\Phi\Sigma}$ are taken to be free values of physical parameters.

III. DARK MATTER

We have two DM candidates which are the lightest fermion F_1 and the lightest mass eigenstate of the scalars $(\eta, \chi_0, \chi'_0)_R$. These DM candidates can be stabilized by the accidental \mathbb{Z}_2 symmetry but not a remnant symmetry of the global $U(1)$ symmetry. This accidental \mathbb{Z}_2 symmetry could be understood as a kind of the accidental symmetry which has been discussed for gauged $U(1)_{B-L}$ in Ref. [84].

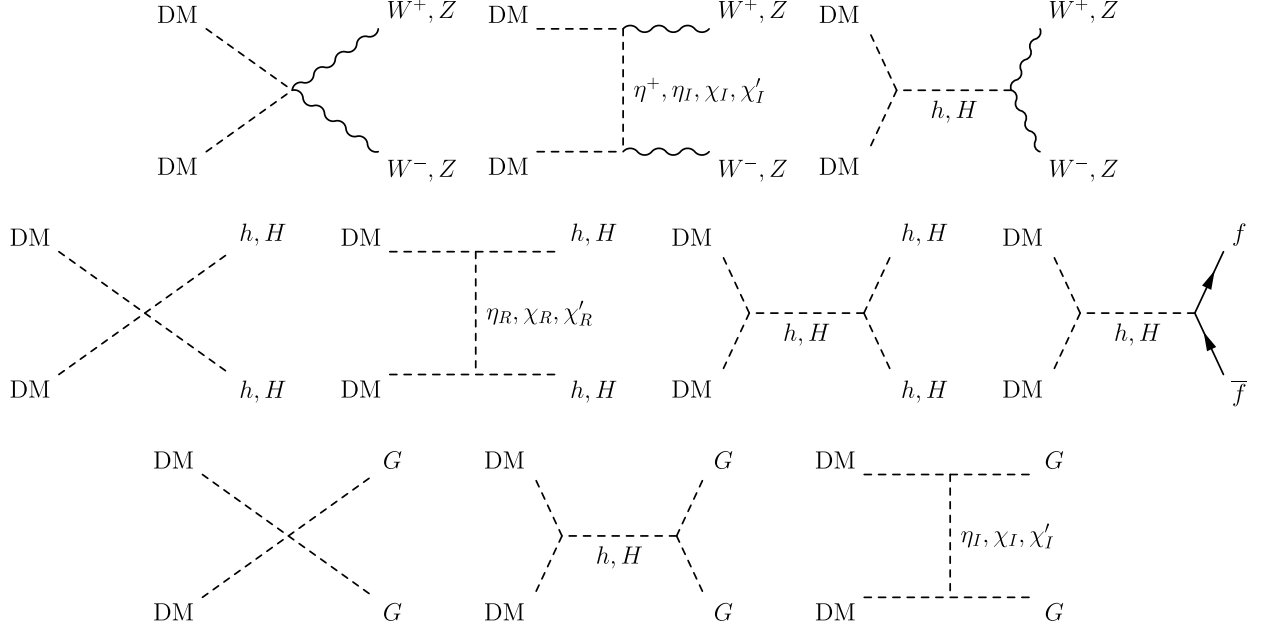


FIG. 3: Diagrams of the DM annihilations where DM is identified as the lightest scalar mass eigenstate.

For the fermionic DM candidate F_1 , the relevant coupling for DM annihilations is only the Yukawa coupling y_η , and the required strength of the Yukawa coupling is $\mathcal{O}(1)$ for the DM mass above the electroweak scale in order to accommodate the observed DM relic density. On the other hand, small coupling $y_\eta \lesssim \mathcal{O}(0.01)$ is needed to evade the LFV constraint as has been discussed in the previous section. Thus the fermionic DM candidate F_1 conflicts with the LFV constraint.⁵

Therefore we identify the lightest mass eigenstate of the scalars as a DM candidate. The mixing angles among $(\eta, \chi_0, \chi'_0)_R$ are induced by the scalar couplings $\lambda, \lambda', \lambda''$, and the magnitude of the mixing angles should roughly be $O_{\text{mix}} \sim 10^{-2}$ in order to reproduce the measured neutrino masses without conflict with the $\mu \rightarrow e\gamma$ process as discussed in the previous section. This order of the magnitude of the mixing angles can be achieved with the parameter setting given in Sec. II A. Since the full scalar potential given by Eq. (II.2) is rather complicated, we take into account only $\lambda, \lambda', \lambda'', \lambda_\Phi, \lambda_{\Phi\Sigma}, \lambda_\Sigma$ and $\lambda_{\Phi\eta}$ for simplicity. Since the required order of the magnitude of the couplings $\lambda, \lambda', \lambda'', \lambda_{\Phi\Sigma}, \lambda_\Sigma$ is very small, they would not affect to the computation of the DM relic density and detection probability. However the coupling $\lambda_{\Phi\Sigma}$ is important to induce the mixing angle $\sin \alpha$, and $\lambda_{\Phi\eta}$ is responsible for direct detection of DM since this coupling generates the dominant contribution to the elastic scattering with nuclei mediated by the SM-like Higgs boson

⁵ Although the LFV constraint may be satisfied by considering a diagonal Yukawa matrix or specific flavor structure, we do not discuss this case.

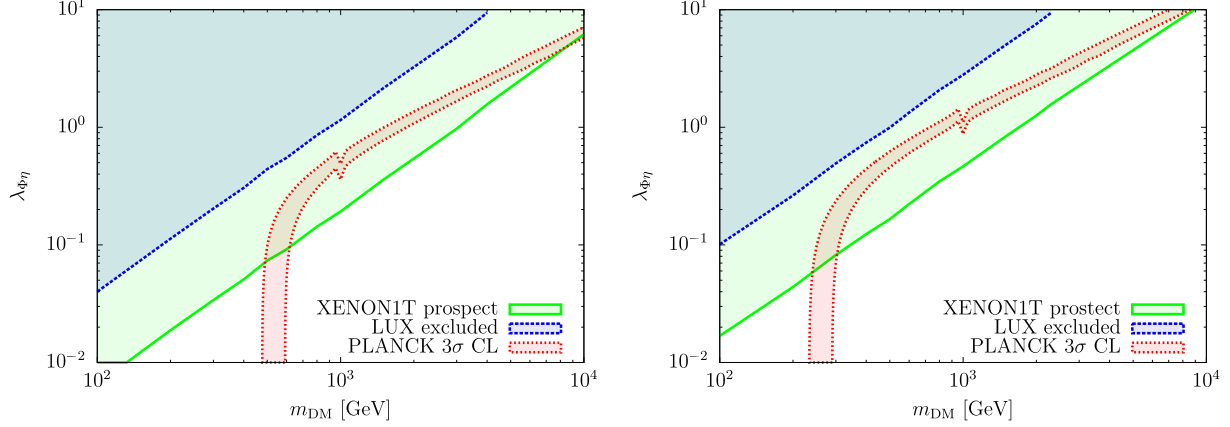


FIG. 4: Numerical results in the $(\lambda_{\Phi\eta}, m_{\text{DM}})$ plane for $(\eta^0, \chi^0, \chi^{0'})_R$ non-degenerate case (left plot) and degenerate case (right plot), where the parameters are fixed to be $m_H = 2$ TeV and $\sin \alpha = 0.1$.

h .

The diagrams of the DM annihilations are shown in Fig. 3. The DM couplings in the scalar potential are basically weak in our parameter setting. However since the scalar DM candidate includes the $SU(2)_L$ doublet inert scalar η_R , DM can annihilate into the gauge bosons via the gauge interactions in order to satisfy the observed DM relic density if the inert doublet scalar component of the DM candidate is sufficiently large. This can be achieved with a smaller $(M_R^2)_{11}$ compared to $(M_R^2)_{22}$ and $(M_R^2)_{33}$ in Eq. (II.6), and we consider such a case. In this case, the annihilation channels in the first line in Fig. 3 become dominant processes to determine the DM relic density.

The DM relic density can be evaluated by using micrOMEGAs [85] and the results are shown in the $(\lambda_{\Phi\eta}, m_{\text{DM}})$ plane in Fig. 4 where the heavier CP-even Higgs boson mass is taken to be $m_H = 2$ TeV and the mixing angle is $\sin \alpha = 0.1$ as an example. Here we take the negative values of the scalar couplings $\lambda, \lambda', \lambda''$. If the scalar couplings are positive, the lightest scalar in the imaginary components becomes DM candidate instead of the real components. The left plot shows the case that the masses of the other two heavier mass eigenstates are twice of the DM mass (lightest state), and the right plot shows the case that the heavier states are degenerate with the DM state. The red colored band represents the region satisfying the observed DM relic density by PLANCK within 3σ confidence level [86]. The blue region is excluded by the direct detection experiment LUX [87], and the green region is expected to be tested by the future direct detection experiment XENON1T [88]. From these plots, one can see that when the lightest DM state is non-degenerate with the other heavier states, DM is close to the inert doublet DM (left plot), because

there is the mass threshold $m_{\text{DM}} \approx 530$ GeV for the left plot in Fig. 4, which is the same property of the inert doublet DM [89]. As well-known, the inert doublet scalar DM candidate can satisfy the observed relic density in the mass ranges of $m_{\text{DM}} \sim 60$ GeV and $m_{\text{DM}} \gtrsim 530$ GeV. On the other hand, the mass threshold can be lower as $m_{\text{DM}} \sim 250$ GeV if χ^0 and $\chi^{0'}$ are degenerate with DM as one can see from the right plot in Fig. 4. This is because the interactions of the singlets χ^0 and $\chi^{0'}$ are described by the scalar potential, and extremely limited in the case of the simplified potential. There is a small resonance feature at $m_{\text{DM}} \sim 1$ TeV due to the channel $\text{DMDM} \rightarrow H^* \rightarrow \text{SMSM}$, however the resonance is not strong because of the small mixing angle $\sin \alpha = 0.1$.

IV. RESONANT LEPTOGENESIS

We consider the thermal leptogenesis in this model [90]. The lepton number asymmetry is expected to be generated through the out-of-equilibrium decay of the lightest Majorana fermions N_1 and X_1 , if we impose the lepton number of F as -1 . Although the Yukawa coupling y_η is required to be smaller than $\mathcal{O}(0.01)$ from the LFV constraint, this is large enough that F and the SM leptons are in the thermal equilibrium. Thus, the generated lepton number asymmetry in the F sector can instantaneously be converted into the SM leptons, and then the baryon number asymmetry can be generated through sphaleron process.

After the global $U(1)$ symmetry breaking, the Yukawa interactions for Majorana fermions are written as

$$\begin{aligned} \mathcal{L} \supset & \bar{F}^c(Y_{N_i}P_L + Y'_{N_i}P_R)N'_i\chi^* + \bar{F}^c(Y_{X_i}P_L + Y'_{X_i}P_R)X'_i\chi^* \\ & + \bar{F}^c(Y_{N_i}P_L + Y'_{N_i}P_R)N'_i\chi' + \bar{F}^c(Y_{X_i}P_L + Y'_{X_i}P_R)X'_i\chi' \\ & + \bar{F}(Y'_{N_i}P_L + Y_{N_i}P_R)N'_i\chi'^* + \bar{F}(Y'_{X_i}P_L + Y_{X_i}P_R)X'_i\chi'^* \\ & + \bar{F}(Y'_{N_i}P_L + Y_{N_i}P_R)N'_i\chi + \bar{F}(Y'_{X_i}P_L + Y_{X_i}P_R)X'_i\chi, \end{aligned} \quad (\text{IV.1})$$

where N'_i and X'_i are expressed as the mass eigenstates of each Majorana fermion. Hereafter, we abbreviate them to N_i and X_i for convenience. The Yukawa couplings are redefined as

$$\begin{aligned} Y_{N_i\chi} &= y_{N_i\chi}^* \cos \theta_i, \quad Y'_{N_i\chi} = y'_{N_i\chi} \sin \theta_i, \\ Y_{X_i\chi} &= -y_{N_i\chi}^* \sin \theta_i, \quad Y'_{X_i\chi} = y'_{X_i\chi} \cos \theta_i, \\ Y_{N_i\chi'} &= y_{N_i\chi'}^* \sin \theta_i, \quad Y'_{N_i\chi'} = y'_{N_i\chi'} \cos \theta_i, \\ Y_{X_i\chi'} &= y_{X_i\chi'}^* \cos \theta_i, \quad Y'_{X_i\chi'} = -y'_{X_i\chi'} \sin \theta_i, \end{aligned} \quad (\text{IV.2})$$

where θ_i is the mixing angle of the i -th generation.

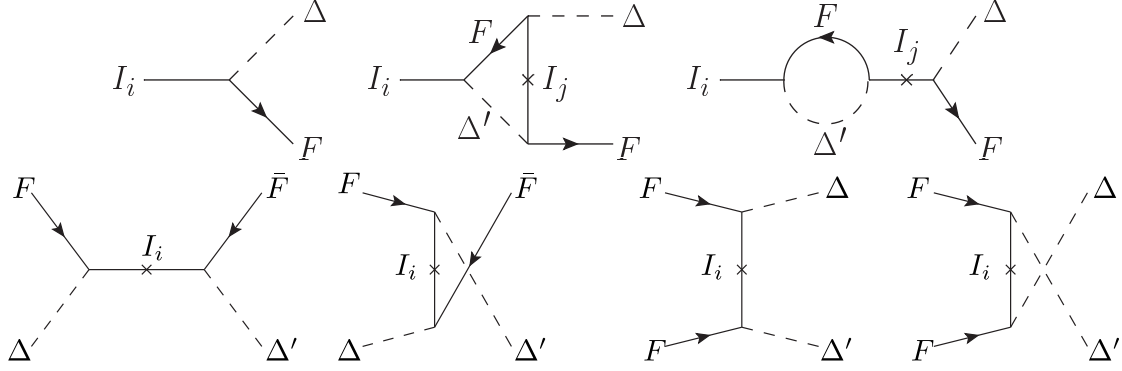


FIG. 5: Lepton number violating decay and scattering processes for leptogenesis where Δ and Δ' represent the new scalar singlets, χ^0 or $\chi^{0'}$ and I_i ($i = 1 - 2$) is N_i or X_i

In this model, we consider that the TeV scale masses of Majorana fermions and $v' \sim 10^7$ GeV so that the annihilation channels $N_1 N_1, X_1 X_1 \rightarrow GG$ are decoupled from the thermal bath before the temperature of the universe $T \sim 10$ TeV. Otherwise the generated lepton asymmetry would be washed out by these lepton number violating processes.

Although the TeV scale of the masses seems to be too small to realize the sufficient CP asymmetry associated with the decay processes for the generation of the required baryon number asymmetry, the generated baryon asymmetry can be enhanced by the resonance effect as known in resonant leptogenesis [91–99]. We define the parameter ϵ as the amplitude of the CP asymmetry. The dominant contribution for the CP asymmetry comes from the interference between the tree diagram and the one-loop self-energy diagram as depicted in the upper line in Fig. 5 and its formula is given by

$$\epsilon \propto \frac{(M_{I_1}^2 - M_{I_2}^2) M_{I_2} \Gamma_{I_2}}{(M_{I_1}^2 - M_{I_2}^2)^2 + M_{I_2}^2 \Gamma_{I_2}^2}, \quad (\text{IV.3})$$

where M_{I_i} and Γ_{I_i} are the mass and the decay rate of I_i ($I = N$ or X) respectively. From this equation, the maximum enhancement is caused when $M_{I_1}^2 - M_{I_2}^2 = M_{I_2} \Gamma_{I_2}$. In this model, we can take a larger Γ_{I_2} compared to that of the second lightest right-handed neutrino in the canonical resonant leptogenesis at TeV scale, since the Yukawa couplings can be large without conflicting with the observed neutrino masses due to the loop suppression. Thus, the required magnitude of the degeneracy of the Majorana fermion mass can be quite milder to generate the sufficient baryon number asymmetry than those in the usual resonant leptogenesis at TeV scale [73, 100]. One may think that the baryon asymmetry should be correlated with the VEV of the singlet scalar Σ since

the $B - L$ breaking occurs with only v' . Indeed this $B - L$ breaking effect is included in the total mass matrix of N_i , X_i and the active neutrinos ν_i . For example, the $B - L$ violating Dirac mass term between ν_i and X_j is induced at one-loop level as can be seen from the left diagram in Fig. 1. Thus the effect of the $B - L$ breaking is included in the masses of N_i and X_i , and one can understand that the baryon asymmetry is generated through the breaking effect in the mass matrix.

We need to take into account washout effects to evaluate the baryon number asymmetry. The generated lepton number asymmetry could be washed out through the lepton number violating 2-2 scattering processes and the inverse decay of I_i . However, if the relevant Yukawa couplings are small enough, these processes can be nearly decoupled before the temperature of the thermal plasma decreases to $T \lesssim M_{I_1}$. Thus, the washout of the generated lepton number asymmetry is expected to be suppressed sufficiently in this period. In order to examine this quantitatively, we numerically solve the coupled Boltzmann equations for the number density of N_1 , X_1 and the lepton number asymmetry. We introduce the number density of N_1 , X_1 and the lepton number asymmetry in the comoving volume as $Y_{N_1} = n_{N_1}/s$, $Y_{X_1} = n_{X_1}/s$ and $Y_F = (n_F - n_{\bar{F}})/s$ respectively, by using the entropy density s and the number densities which are expressed by n_{N_1} , n_{X_1} , n_F and $n_{\bar{F}}$. Their equilibrium values are given by $Y_{I_1}^{\text{eq}} = \frac{45}{2\pi^4 g_*} z^2 K_2(z)$, where z is defined by $z = M_{I_1}/T$, g_* is the number of relativistic degrees of freedom and $K_2(z)$ is the modified Bessel function of the second kind with the order 2. Since we assume Y_F is immediately translated into the SM leptons as we mentioned above, we use the relation $B = \frac{8}{23}(B - L)$ which is derived from the chemical equilibrium condition in this model, and the baryon number asymmetry Y_B in the present Universe is estimated as $Y_B = -\frac{8}{23}Y_F(z_{\text{EW}})$, where $z_{\text{EW}} = M_{I_1}/T_{\text{EW}}$ is related to the sphaleron decoupling temperature T_{EW} .

The coupled Boltzmann equations for the leptogenesis in our model are written as [101]

$$\frac{dY_{N_1}}{dz} = -\frac{z}{sH} \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) \gamma_{N_1}^D, \quad (\text{IV.4})$$

$$\frac{dY_{X_1}}{dz} = -\frac{z}{sH} \left(\frac{Y_{X_1}}{Y_{X_1}^{\text{eq}}} - 1 \right) \gamma_{X_1}^D, \quad (\text{IV.5})$$

$$\begin{aligned} \frac{dY_F}{dz} = \frac{z}{sH} \left\{ \epsilon_N \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) \gamma_{N_1}^D + \epsilon_X \left(\frac{Y_{X_1}}{Y_{X_1}^{\text{eq}}} - 1 \right) \gamma_{X_1}^D \right. \\ \left. - \frac{2Y_F}{Y_F^{\text{eq}}} \left[\sum_i \frac{\gamma_{N_i} + \gamma_{X_i}}{4} + \sum_{\Delta, \Delta'} (\gamma_{F\Delta F\Delta'} + \gamma_{FF\Delta\Delta'}) \right] \right\}, \end{aligned} \quad (\text{IV.6})$$

where $\gamma_{I_i}^D$ is defined by

$$\gamma_{I_i}^D = \frac{M_{I_i}^2 T}{\pi^2} K_1\left(\frac{M_{I_i}}{T}\right) \Gamma_{I_i}^D, \quad (\text{IV.7})$$

with the modified Bessel function of the second kind $K_1(z)$ with the order 1, γ_{abij} is the reaction density for the scattering process $ab \leftrightarrow ij$ which is given by

$$\gamma_{abij} = \frac{T}{64\pi^4} \int_{s_{\min}}^{\infty} ds \, \hat{\sigma}_{abij}(s) \sqrt{s} K_1\left(\frac{\sqrt{s}}{T}\right), \quad (\text{IV.8})$$

with $s_{\min} = \max[(m_a + m_b)^2, (m_i + m_j)^2]$ and the reduced cross section $\hat{\sigma}_{abij}(s)$. There are two kinds of the processes $F\Delta \leftrightarrow \bar{F}\Delta'$ and $FF \leftrightarrow \Delta\Delta'$ for the scattering processes as depicted in the bottom line of Fig. 5. The reduced cross section in our model is rather complicated since a lot of particles exist, but can be straightforwardly computed from the Lagrangian as same as Ref. [102, 103]. We solve the coupled Boltzmann equations numerically.

The decay of the lightest Majorana fermions should be out of thermal equilibrium so that the lepton number asymmetry can be generated through their decays. If we express the Hubble parameter as H , this condition is given by $H > \Gamma_{I_1}$ at $T \sim M_{I_1}$. Since we assume that each of the Majorana fermion mass is $M_{I_1} = 5$ TeV, $M_{I_2} = M_{I_1}(1 + \delta M)$ and $M_{I_3} = 6$ TeV, the Yukawa couplings of N_1 and X_1 should be $\mathcal{O}(10^{-8})$. On the other hand, the rest of the Yukawa couplings should be $\mathcal{O}(10^{-2})$ in order to generate the appropriate neutrino masses. Here we set Yukawa couplings in Eq. (IV.2) to be $y_{I_1\chi} = 1.5 \times 10^{-8}$ and $y_{I_{2,3}\chi} = 10^{-2}$ and $\theta_i = \frac{\pi}{4}$. We show the result for $\delta M = 10^{-3}$ in Fig. 6 as an example, and also the generated baryon number asymmetry for each value of δM and M_{I_1} in Fig. 7. Through this analysis, the masses of F , χ and χ' are fixed to be $M_F = 1.5$ TeV and $M_\chi = M_{\chi'} = 1.2$ TeV respectively. These parameter set satisfies the condition for the DM phenomenology we discussed in the previous section.

From the left panel in Fig. 6, we can see that the required baryon number asymmetry Y_B can be obtained in this model. In the right panel, we plot the behavior of the relevant reaction rate for each process. This panel shows that the reaction rates of the inverse decay process and the lepton number-violating process induced by the s-channel I_i exchange are quite large for a long time. Thus the baryon number asymmetry cannot be generated quickly until rather a late period. After $T \sim M_{I_1}$, the generated baryon number asymmetry gradually increases and then the required value can be realized. This is because these processes are suppressed by the Boltzmann factor.

We show the relation between the generated baryon number asymmetry and the mass degeneracy of Majorana fermions in Fig. 7. Notice here that the generated baryon number asymmetry is always smaller than the required value in the case $M_{I_1} \sim 3$ TeV. In this model, we can realize

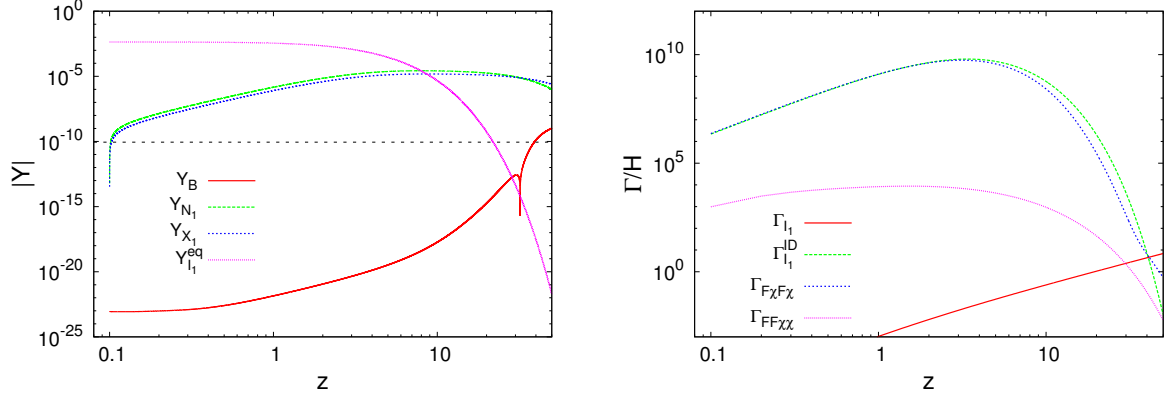


FIG. 6: Left panel: Evolution of Y_B , Y_{N_1} and Y_{X_1} where the horizontal black dashed line represents the required value of the baryon number asymmetry. Right panel: Reaction rates Γ/H of the processes that have crucial effects for the baryon number asymmetry.

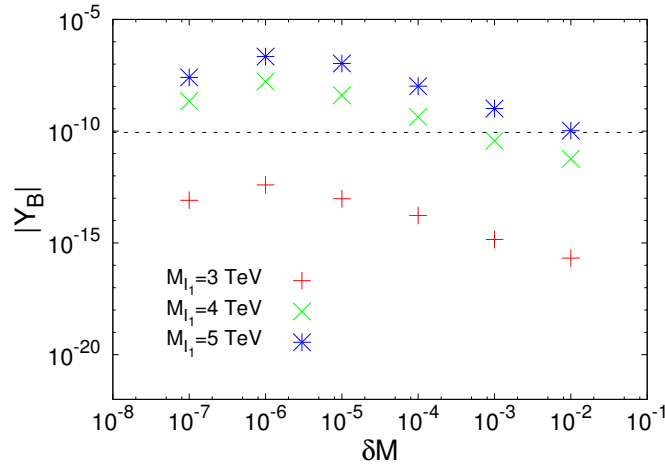


FIG. 7: δM dependence of the generated baryon number asymmetry Y_B with the different value of M_{I_1} . The value of M_{I_1} is fixed to 3 TeV, 4 TeV and 5 TeV respectively.

the large Yukawa couplings to explain the small neutrino mass due to the two-loop effects and then Γ_{I_2} becomes larger compared to tree and one-loop neutrino mass models. Thus, the required mass degeneracy can be milder. However, the large Yukawa couplings cause the large washout effects. Since the Boltzmann suppression does not work well in the case of small M_{I_1} , the large washout effects remain until quite a late period compared to the heavier cases. Thus the most of the generated baryon number asymmetry is washed out. We find that the observed baryon number asymmetry can be generated when the mass degeneracy is $\delta M = (10^{-3} - 10^{-2})$ for $M_{I_1} = 4 - 5$ TeV, as can be seen in Fig. 7. As we mentioned above, the magnitude of this mass degeneracy is

quite milder than the canonical seesaw case for each value of M_{I_1} due to the loop effects.⁶

V. CONCLUSIONS

We have studied a two-loop induced radiative neutrino model at TeV scale with a $U(1)$ global symmetry, in which two types of DM candidates (the lightest one of fermion or scalar) can be involved. The loop-induced neutrino masses have been evaluated appropriately and phenomenology of DM has also been discussed. The fermionic DM candidate is disfavored if we reproduce the measured neutrino masses and take into account the constraint from LFV with the same order of all the elements of y_η . Then we have found that the scalar can be a good DM candidate which satisfies the observed relic density and the DM direct detection bound. We also found that the direct detection rate of DM is controlled by the coupling $\lambda_{\Phi\eta}$ and some parameter region can be testable by the next future direct detection experiment XENON1T.

We have discussed baryon number asymmetry through the resonant leptogenesis with multi-sources scenario, in which the large Yukawa couplings (that is required to compensate the tiny neutrino masses at two-loop level) make the large CP asymmetry, but also cause the large washout effects. We have shown that the required baryon number asymmetry can be obtained for the parameter region *i.e.*, $\delta M = (10^{-3} - 10^{-2})$ for $M_{I_1} = 4 - 5$ TeV, where $I \equiv N$ or X . In this case, the lightest Majorana fermions should satisfy $M_{I_1} \gtrsim 3$ TeV to suppress the washout effects by the Boltzmann factor. For larger M_{I_1} , the required magnitude of the mass degeneracy is rather milder than the canonical seesaw case even at TeV scale models.

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⁶ The typical scale of δM is $10^{-10} - 10^{-8}$. It is also worth mentioning that the Ma model has an solution for $\delta M \leq 10^{-6.5}$ [73].

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